

Constraints on 750 GeV colorless Q -onia from running couplings

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Abstract

We present yet another composite model explaining the relatively broad peak in the diphoton invariant mass distribution around 750 GeV recently observed at the LHC experiments. We consider the excess originates in bound states of vector-like fermions Q transforming under the electroweak group (but not QCD) of the standard model and which are also charged under a new $SU(N_{TC})$ QCD-like confining force. Since the new uncolored fields transform as $SU(2)$ doublets they can acquire a mass through Yukawa interactions with the electroweak Higgs doublet, as quarks and leptons. We parametrize the Q -onium bound state using the Coulomb approximation and give a numerical fit to the diphoton excess consistent with no strong coupling effects up to about 30 TeV, which is the energy scale for next generation colliders. The range of allowed hypercharge Y_Q is given by $2.26 \leq Y_Q \leq 2.53$ for $N_{TC} = 2$ and $2.17 \leq Y_Q \leq 2.31$ for $N_{TC} = 3$. The smoking gun for the model would be the discovery of Q -onia decaying into Higgs and Z -boson.

I. INTRODUCTION

In December 2015, ATLAS [1] and CMS [2] famously announced the observation of a peak in the diphoton mass distribution around 750 GeV, using (respectively) 3.2 fb^{-1} and 2.6 fb^{-1} of data recorded at a center-of-mass energy $\sqrt{s} = 13 \text{ TeV}$. The diphoton excesses could be interpreted as the decay products of a new massive particle X , with spin 0, 2, or higher [3]. Assuming a narrow width approximation ATLAS gives a local significance of 3.6σ , or else a global significance of 2.0σ when the look-elsewhere-effect in the mass range $M_X/\text{GeV} \in [200 - 2000]$ is accounted for. Signal-plus-background fits were also implemented for a broad signal component with a large decay width. The largest deviation from the background-only hypothesis corresponds to $M_X \sim 750 \text{ GeV}$ with a total width $\Gamma_{\text{total}} \sim 45 \text{ GeV}$. The local and global significances evaluated for the broad resonance fit are roughly 0.3 higher than that for the fit using the narrow width approximation, corresponding to 3.9σ and 2.3σ , respectively. The CMS data gives a local significance of 2.6σ and a global significance smaller than 1.2σ . Fitting the LHC data at $\sqrt{s} = 13 \text{ TeV}$ to a resonant structure leads to a cross section times branching fraction of

$$\sigma_{\text{LHC13}}(pp \rightarrow X + \text{anything}) \times \mathcal{B}(X \rightarrow \gamma\gamma) \sim \begin{cases} (10 \pm 3) \text{ fb} & \text{ATLAS} \\ (6 \pm 3) \text{ fb} & \text{CMS} \end{cases}, \quad (1)$$

at 1σ [5]. However, the LHC data at $\sqrt{s} = 8 \text{ TeV}$ show no excess over standard model processes (SM) [6, 7], placing a critical upper bound on the cross section times branching fraction: $\sigma_{\text{LHC8}}(pp \rightarrow X + \text{anything}) \times \mathcal{B}(X \rightarrow \gamma\gamma) < 2.00 \text{ fb}$ at 95% CL [7]. Altogether, the LHC8 data would be compatible with the LHC13 data within $\sim 2\sigma$, if the diphoton cross section grows by more than about a factor of 3 or 3.5.

Quite recently, ATLAS and CMS updated their diphoton resonance searches [8–10]. ATLAS reanalyzed the 3.2 fb^{-1} of data, targeting separately spin-0 and spin-2 resonances. For spin-0, the most significant deviation from the background-only hypothesis corresponds to $M_X \sim 750 \text{ GeV}$ and $\Gamma_{\text{total}} \sim 45 \text{ GeV}$. The local significance is now increased to 3.9σ but the global significance remains at the 2σ level. For the spin-2 resonance, both the local and global significances are reduced down to 3.6σ and 1.8σ , respectively. The new CMS analysis includes additional data (recorded in 2015 while the magnet was not operated) for a total of 3.3 fb^{-1} . The largest excess is observed for $M_X = 760 \text{ GeV}$ and $\Gamma_{\text{total}} \approx 11 \text{ GeV}$, and has a local significance of 2.8σ for spin-0 and 2.9σ spin-2 hypothesis. After taking into account the effect of searching for several signal hypotheses, the significance of the excess is reduced to $< 1\sigma$. CMS also communicated a combined search with data recorded at $\sqrt{s} = 13 \text{ TeV}$ and $\sqrt{s} = 8 \text{ TeV}$. For the combined analysis, the largest excess is observed at $M_X = 750 \text{ GeV}$ and $\Gamma_{\text{total}} = 0.1 \text{ GeV}$. The local and global significances are $\approx 3.4\sigma$ and 1.6σ , respectively.

Among many possible interpretations [11], one of the most attractive type of models is that of pairs of new heavy fermions which produce resonant signals of their near-threshold bound states. It seems natural to classify the various models in accordance with the fermion properties. The simplest explanation for the structure underlying the bump in the diphoton spectrum could be that of a QCD bound state of a heavy vector-like quark Q , with a mass around 375 GeV [12–18]. For this model, the signal rate can be explained by a scalar $Q\bar{Q}$ bound state η_Q , with Q transforming as $(3, 1, -4/3)$ under the $SU(3) \otimes SU(2) \otimes U(1)_Y$ gauge groups of the SM [16, 17]. Actually, in analogy to η_c and J/ψ , the QCD bound states would contain both scalars and vectors. We may note in passing that the vector J/ψ -like state will not only decay into $\gamma\gamma$ but also into dilepton topologies. The latter is severely constrained

by ATLAS [19] and CMS [20] direct searches; however, since the J/ψ -like state can only be produced through $q\bar{q}$ annihilation, its production cross section at the LHC is significantly smaller than that of the $Q\bar{Q}$ scalar state.

A somewhat related model interprets the excess of diphoton events as arising also from the decay of a Q -onium state η_Q , but with the colored fermions Q bound by a hidden confining $SU(N)$ [21–29]. One particularly interesting possibility to describe the hidden gauge dynamics is given by $SU(N_{\text{TC}})$, with N_{TF} techni-flavors. This is because the dynamics of this theory is renowned from QCD and can be also understood in the large N_{TC} limit: the gauge theory is asymptotically free provided it satisfies the familiar bound on the number of techni-flavors and confines at a scale Λ_{TC} . To avoid the strict constraints common to old techni-color models the heavy-fermions must be in a vectorial representation of the SM and in the fundamental \mathbf{N}_{TC} of $SU(N_{\text{TC}})$; namely

$$Q = \sum_{i=1}^{N_s} Q_i, \quad \text{with} \quad Q_i = (\mathbf{N}_{\text{TC}}, \mathbf{R}_i) \oplus (\bar{\mathbf{N}}_{\text{TC}}, \bar{\mathbf{R}}_i), \quad (2)$$

where \mathbf{R}_i denotes a generic SM representation and N_s is the number of species with mass below the confinement scale [26]. Yet a third dynamics is possible if the heavy fermions Q are colorless particles, bound by a confining $SU(N_{\text{TC}})$ [30].

At this point it is important to stress two fundamental differences between these three models. Firstly, while η_Q and $\eta_{\bar{Q}}$ could be produced either via gluon or photon fusion in LHC collisions, the production of η_Q could only proceed via photon fusion. Secondly, if the heavy particles belong to the $SU(2)$ doublet and feel the strong $SU(3)$ color interactions, then the 125 GeV Higgs signal strength in the $gg \rightarrow H \rightarrow \gamma\gamma$ channel would be significantly modified. ATLAS [31] and CMS [32] data would then place severe constraints on model parameters unless the coupling of Q or \bar{Q} to the Higgs is small. However, Q transforms as an $SU(3)$ singlet and so it does not contribute to $gg \rightarrow H \rightarrow \gamma\gamma$. Therefore, the Q 's could get mass like quarks and leptons, as $SU(2)$ doublets and singlets can form a mass term after being Higgsed. This is precisely what motivates our study.

In this paper we investigate the phenomenology of Q -onia states of a hidden $SU(N_{\text{TC}})$. We combine the requirements to reproduce the LHC diphoton signal with those arising from the renormalization group (RG) equations to constrain the parameter space of Q fermions transforming under the electroweak group (but not QCD) of the SM. For related RG studies in which the 750 GeV excitation is not a bound state, see [33–37]. Before proceeding we note that the LHC phenomenology of Q -onia states, with constituents that carry only hypercharge, but no SM $SU(3)$ or $SU(2)$ quantum numbers has been presented in [30]. Though at first sight our study seems quite similar to the analysis in [30], it differs in a fundamental aspect: herein we consider colorless fermions that transform as $SU(2)$ doublets and thus the heavy Q particles can obtain a mass, as all other SM fermions, through the Higgs mechanism. The layout of the paper is as follows. In Sec. II we introduce the Lagrangian of the model with a new confining gauge interaction and derive the RG equations. In Sec. III we discuss the LHC phenomenology of our set up. We parametrize the Q -onium bound state using the Coulomb approximation [13] and give a numerical fit to the diphoton excess consistent with the running couplings. We request no strong coupling effects up to about 30 TeV, which is the energy scale for next generation colliders [38]. In Sec. IV we verify consistency with early universe cosmology. Our conclusions are collected in Sec. V.

II. LAGRANGIAN AND RENORMALIZATION GROUP EQUATIONS

To develop our program in the simplest way, we will work within the construct of a minimal model in which we consider one generation of heavy colorless fermions, which contains two $SU(2)$ doublets (one left- and one right-handed to make the representation vectorial, i.e. symmetric under parity) and four chiral fermions that transform as $SU(2)$ singlets (two of them are left-handed and the other two right-handed). We label the four flavors as club, diamond, heart, and spade, and so we named the \mathcal{Q} particles “qards” accordingly.¹ The Lagrangian for the technicolor qards is given by

$$\begin{aligned} \mathcal{L}_{\mathcal{Q}} = & \left(i\overline{\mathcal{Q}}_L \gamma_\mu D^\mu \mathcal{Q}_L + i\overline{\mathcal{Q}}_R^\star \gamma_\mu D^\mu \mathcal{Q}_R^\star + i\overline{\mathcal{Q}}_R^\diamond \gamma_\mu D^\mu \mathcal{Q}_R^\diamond + \right. \\ & \left. + i\overline{\mathcal{Q}}_R \gamma_\mu D^\mu \mathcal{Q}_R + i\overline{\mathcal{Q}}_L^\diamond \gamma_\mu D^\mu \mathcal{Q}_L^\diamond + i\overline{\mathcal{Q}}_L^\star \gamma_\mu D^\mu \mathcal{Q}_L^\star \right) + \text{h.c.}, \end{aligned} \quad (3)$$

where the covariant derivative for the $SU(2)$ doublets $\mathcal{Q}_{L,R}$,

$$\mathcal{Q}_L = \begin{pmatrix} \mathcal{Q}_L^\star \\ \mathcal{Q}_L^\diamond \end{pmatrix} \quad \text{and} \quad \mathcal{Q}_R = \begin{pmatrix} \mathcal{Q}_R^\diamond \\ \mathcal{Q}_R^\star \end{pmatrix}, \quad (4)$$

reads

$$D_\mu = \partial_\mu - ig_2 \tau^a A_\mu^a - ig_{\text{TC}} \tilde{T}^a \tilde{G}_\mu^a - ig_Y Y_{\mathcal{Q}} B_\mu, \quad (5)$$

while those for $SU(2)$ singlets \mathcal{Q}_L^\diamond , \mathcal{Q}_L^\star , \mathcal{Q}_R^\star , \mathcal{Q}_R^\diamond are (plus sign for \mathcal{Q}_R^\star , \mathcal{Q}_L^\diamond)

$$D_\mu = \partial_\mu - ig_{\text{TC}} \tilde{T}^a \tilde{G}_\mu^a - ig_Y \left(Y_{\mathcal{Q}} \pm \frac{1}{2} \right) B_\mu. \quad (6)$$

The Yukawa interactions that provide the masses to the qards are:

$$\mathcal{L}_Y = -\mathcal{Y}_\diamond \overline{\mathcal{Q}}_L \cdot H \mathcal{Q}_R^\diamond - \mathcal{Y}_\star \overline{\mathcal{Q}}_L \cdot \tilde{H} \mathcal{Q}_R^\star - \mathcal{Y}_\star \overline{\mathcal{Q}}_L^\star \mathcal{Q}_R \cdot H - \mathcal{Y}_\diamond \overline{\mathcal{Q}}_L^\diamond \mathcal{Q}_R \cdot \tilde{H} + \text{h.c.} \quad (7)$$

where $\tilde{H} = i\sigma^2 H^*$. The couplings of the technicolor qards with the Higgs are the same as those with SM fermions and hence they get masses in the same way. Substituting in (7) the Higgs vacuum expectation value,

$$\langle H \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v \end{pmatrix}, \quad (8)$$

we have the mass terms

$$\mathcal{L}_m = -\frac{1}{\sqrt{2}} \mathcal{Y}_\diamond \overline{\mathcal{Q}}_L^\diamond \mathcal{Q}_R^\diamond v - \frac{1}{\sqrt{2}} \mathcal{Y}_\star \overline{\mathcal{Q}}_L^\star \mathcal{Q}_R^\star v - \frac{1}{\sqrt{2}} \mathcal{Y}_\star \overline{\mathcal{Q}}_L^\star \mathcal{Q}_R^\star v - \frac{1}{\sqrt{2}} \mathcal{Y}_\diamond \overline{\mathcal{Q}}_L^\diamond \mathcal{Q}_R^\diamond v + \text{h.c.} \quad (9)$$

The RG equation for the coupling of an $SU(N_c)$ gauge field with N_f fundamental matter reads

$$\frac{d}{dt} g(t) = -\frac{g^3}{(4\pi)^2} \left(\frac{11N_c}{3} - \frac{2N_f}{3} \right). \quad (10)$$

¹ Note that qards are quirks transforming under the electroweak group (but not QCD) of the SM [39–43].

Throughout we take both the left and right-handed technicolor $SU(N_{\text{TC}})$ qards to be in the representation of

$$(\mathbf{N}_{\text{TC}}, \mathbf{1}, \mathbf{2} \oplus \mathbf{1} \oplus \mathbf{1})_Y \quad (11)$$

under $SU(N_{\text{TC}}) \otimes SU(3) \otimes SU(2) \times U(1)_Y$, satisfying a relation akin to (2),

$$\mathcal{Q} = (\mathbf{N}_{\text{TC}}, \mathbf{1}, \mathbf{2} \oplus \mathbf{1} \oplus \mathbf{1})_Y \oplus (\bar{\mathbf{N}}_{\text{TC}}, \mathbf{1}, \mathbf{2} \oplus \mathbf{1} \oplus \mathbf{1})_{-Y}. \quad (12)$$

Thus, the RG equation become

$$\frac{d}{dt}g_{\text{TC}} = \frac{g_{\text{TC}}^3}{16\pi^2} \left[-\frac{11}{3}N_{\text{TC}} + \frac{2}{3} \times 4 \right] = \frac{g_{\text{TC}}^3}{16\pi^2} \left(-\frac{11}{3}N_{\text{TC}} + \frac{8}{3} \right), \quad (13)$$

For the $SU(3)$ and $SU(2)$ gauge couplings g_3 and g_2 , the RG equations are

$$\frac{d}{dt}g_3 = \frac{g_3^3}{16\pi^2} \left[-11 + \frac{2}{3}(2 \times 3) \right] = \frac{g_3^3}{16\pi^2} (-7), \quad (14)$$

and

$$\frac{d}{dt}g_2 = \frac{g_2^3}{16\pi^2} \left[-\frac{22}{3} + \frac{2}{3} \left(\frac{3}{2} \times 3 + \frac{1}{2} \times 3 + N_{\text{TC}} \right) + \frac{1}{6} \right] = \frac{g_2^3}{16\pi^2} \left(-\frac{19}{6} + \frac{2}{3}N_{\text{TC}} \right). \quad (15)$$

The RG running of the hypercharge coupling can be worked out in a similar fashion,

$$\frac{d}{dt}g_Y = \frac{1}{16\pi^2} \left\{ \frac{2}{3} \sum_f Y_f^2 + \frac{1}{3} \sum_s Y_s^2 \right\} g_Y^3, \quad (16)$$

where the sum is over the hypercharges of chiral fermions f and complex scalars s . For the case at hand (16) can be rewritten as

$$\frac{dg_Y}{dt} = \frac{g_Y^3}{16\pi^2} \left\{ \frac{2}{3} \times 2 \times \left[2N_{\text{TC}}Y_{\mathcal{Q}}^2 + N_{\text{TC}} \left(Y_{\mathcal{Q}} - \frac{1}{2} \right)^2 + N_{\text{TC}} \left(Y_{\mathcal{Q}} + \frac{1}{2} \right)^2 \right] + \frac{41}{6} \right\}, \quad (17)$$

where $Y_{\mathcal{Q}}$ is the hypercharge of the technicolor qards in the $SU(2)$ doublet. Notice that the qards in the doublet of $SU(2)$ have charge $Y_{\mathcal{Q}}$ while those in the singlet have charges $Y_{\mathcal{Q}} \pm \frac{1}{2}$ in order to allow Yukawa coupling. The extra factor of 2 follows from the fact that we have doublets for both left- and right-handed qards.

We denote the Yukawa coupling of the qard with hypercharge $Y_{\mathcal{Q}} \pm \frac{1}{2}$ as \mathcal{Y}_{\pm} . The RG equation for \mathcal{Y}_{\pm} is given by

$$\frac{d}{dt}\mathcal{Y}_{\pm} = \frac{\mathcal{Y}_{\pm}}{16\pi^2} \left[\frac{2N_{\text{TC}} + 3}{2} \mathcal{Y}_{\pm}^2 + \frac{3}{2} \mathcal{Y}_{\mp}^2 - \frac{3(N_{\text{TC}}^2 - 1)}{N_{\text{TC}}} g_{\text{TC}}^2 - \frac{9}{4} g_2^2 - \left(\frac{3}{4} \pm 3Y_{\mathcal{Q}} + 6Y_{\mathcal{Q}}^2 \right) g_Y^2 \right]. \quad (18)$$

The $SU(2)$ quantum number is the same as in SM and hence so is the contribution from $SU(2)$ gauge fields. The only difference is in the hypercharges. It can be easily checked that for $Y_{\mathcal{Q}} = 1/6$, the g_Y contribution $(3/4 + 3Y_{\mathcal{Q}} + 6Y_{\mathcal{Q}}^2)$ gives $17/12$, which is the value for the SM.

III. PLAYING QARDS AT THE LHC

For simplicity, we further assume the initial value of $\mathcal{Y}_+ = \mathcal{Y}_-$, and so there are four possible \mathcal{Q} -onium bound states that can contribute the diphoton excess

$$\bar{\mathcal{Q}}^\heartsuit \mathcal{Q}^\heartsuit, \quad \bar{\mathcal{Q}}^\spadesuit \mathcal{Q}^\spadesuit, \quad \bar{\mathcal{Q}}^\clubsuit \mathcal{Q}^\clubsuit, \quad \bar{\mathcal{Q}}^\diamond \mathcal{Q}^\diamond, \quad (19)$$

where we have already combined the left and right techni-qards to form Dirac spinors. All of them are spin-0 ($\eta_{\mathcal{Q}}$) and have equal masses $M_X \sim 2m_{\mathcal{Q}}$.

Following [12, 13] we describe the $SU(N_{\text{TC}})$ binding potential in the Coulomb approximation. The radial wave function $R^{(n,l)}(r)$ follows from that of the hydrogen atom

$$\left(\frac{|R^{(n,l)}(0)|^2}{M_X^3} \right)_{\text{Coul}} = \frac{1}{16n^3} (C_N \bar{\alpha}_{\text{TC}})^3, \quad (20)$$

where n is the principal quantum number, l is the orbital angular momentum, and $\bar{\alpha}_{\text{TC}} \equiv \alpha_{\text{TC}}(a_0^{-1})$ is the techni-color gauge coupling in the $\overline{\text{MS}}$ scheme, with a_0 the Bohr radius of the bound state [30]. For fermions in the fundamental representation, $C_N = (N_{\text{TC}}^2 - 1)/(2N_{\text{TC}})$. The Coulomb approximation is reliable when the non-perturbative effect is small. In what follows, we demand that the inverse Bohr radius, which is the characteristic scale of the bound state dynamics, is above the confinement scale:

$$a_0^{-1} \sim \alpha_{\text{TC}}(a_0^{-1}) M_X > \Lambda_{\text{TC}} \quad (21)$$

where the confinement scale Λ_{TC} reads

$$\Lambda_{\text{TC}} \sim M_X \exp \left[-\frac{6\pi}{(11N_{\text{TC}} - 2N_f)\alpha_{\text{TC}}(M_X)} \right]. \quad (22)$$

Intuitively, (21) states that the perturbative treatment (like the Coulomb approximation) breaks down below the confinement scale. The value of a_0^{-1} ranges from 58 GeV to about 320 GeV depending on the value of N_{TC} and $Y_{\mathcal{Q}}$.

The four resonance states can be produced via photon fusion [44–50]. The total photo-production cross section at LHC13 can be parametrized by

$$\sigma_{\text{LHC13}}(\gamma\gamma \rightarrow \eta_{\mathcal{Q}} \rightarrow \gamma\gamma) = 4.1 \text{ pb} \left(\frac{\Gamma_{\text{total}}}{45 \text{ GeV}} \right) \mathcal{B}^2(\eta_{\mathcal{Q}} \rightarrow \gamma\gamma) \quad (23)$$

and the one at LHC8 by

$$\sigma_{\text{LHC8}}(\gamma\gamma \rightarrow \eta_{\mathcal{Q}} \rightarrow \gamma\gamma) = 1.4 \text{ pb} \left(\frac{\Gamma_{\text{total}}}{45 \text{ GeV}} \right) \mathcal{B}^2(\eta_{\mathcal{Q}} \rightarrow \gamma\gamma), \quad (24)$$

showing consistency with the 95% CL upper limit [7]. Actually, the ratio of the LHC13/LHC8 partonic luminosity is largely dominated by systematic uncertainties driven by the parton distribution functions. The luminosity ratio is [48]

$$\frac{\mathcal{L}_{\gamma\gamma}(\sqrt{s} = 13 \text{ TeV})}{\mathcal{L}_{\gamma\gamma}(\sqrt{s} = 8 \text{ TeV})} = 3_{-0.2}^{+0.1}, \quad 2.65 \pm 0.15, \quad 2.1 \pm 0.4, \quad (25)$$

for CT14QED [51], MRST2004 [52], and NNPDF2.3 [53]; respectively. We note that the predictions of NNPDF2.3 are only marginally compatible with LHC8 data [7].

The decay width of η_Q to diphotons is given by [30]

$$\frac{\Gamma(\eta_Q \rightarrow \gamma\gamma)}{M_X} = 4N_{\text{TC}} \left(Y_Q + \frac{1}{2}\right)^4 \alpha^2 \frac{|R(0)|^2}{M_X^3}, \quad (26)$$

where $R(0) \equiv R^{(1,0)}(0)$. The total width is dominated by two major channels. One of them is to techni-gluons $\Gamma(\eta_Q \rightarrow g_{\text{TC}}g_{\text{TC}})$, with

$$\frac{\Gamma(\eta_Q \rightarrow g_{\text{TC}}g_{\text{TC}})}{\Gamma(\eta_Q \rightarrow \gamma\gamma)} = \frac{(N_{\text{TC}}^2 - 1)}{4N_{\text{TC}}^2} \frac{\alpha_{\text{TC}}^2}{(Y_Q + \frac{1}{2})^4 \alpha^2}. \quad (27)$$

and the other is $\Gamma(\eta_Q \rightarrow ZH)$ [12]

$$\frac{\Gamma(\eta_Q \rightarrow ZH)}{\Gamma(\eta_Q \rightarrow \gamma\gamma)} = \frac{M_X^4}{4M_Z^4} \frac{a^2 \alpha_Z^2 [(1 - R_H - R_Z)^2 - 4R_H R_Z]^{3/2}}{4(Y_Q + \frac{1}{2})^4 \alpha^2}, \quad (28)$$

where $R_i = (M_i/M_X)^2$ and M_i ($i = H, Z$) are the masses of Higgs and Z -boson. Also $\alpha_Z = \alpha/(\sin^2 \theta_W \cos^2 \theta_W)$ and $a = \frac{1}{2}(I_{3L} - I_{3R})$, where I_{3L} (I_{3R}) is the isospin for the left- and right-handed techni-quarks. Notice that the decay width into ZH is enhanced by a factor of $\frac{M_X^4}{M_Z^4}$ due to the longitudinal mode of the Z -boson. The decay width into HH would receive similar enhancement had it not been forbidden by CP symmetry (HH has $J^{\text{PC}} = 0^{++}$ while η_Q has $J^{\text{PC}} = 0^{-+}$). We will determine α_{TC} (as a function of N_{TC}, Y_Q) by fitting the predicted production cross section times branching to the observed value. Before proceeding we note that the ATLAS excess is quite broad and probably with a large uncertainty. The CMS excess, however, is smaller and has no clear preference for a large width. This seems to indicate that the ATLAS excess could be a real signal combined with a large fluctuation, making the excess appear larger and wider than the underlying physical signal. Throughout we assume the resonance needs to have a signal [17]

$$\sigma_{\text{LHC13}}(pp \rightarrow X + \text{anything}) \times \mathcal{B}(X \rightarrow \gamma\gamma) \approx 3 - 6 \text{ fb}. \quad (29)$$

Note that the 2 bound states of electric charge $Y_Q + \frac{1}{2}$ contribute equally to the total cross section and dwarf the contributions from the other two states of charge $Y_Q - \frac{1}{2}$ because σ is proportional to the 8th power of the electric charge.

Our results are encapsulated in Fig. 1, where we show the region of the parameter space that can explain the diphoton signal and satisfy the bound $\Gamma(\eta_Q \rightarrow ZH) < 10 \Gamma(\eta_Q \rightarrow \gamma\gamma)$. The (blue) banana-shape region in the top panel is obtained by demanding both that the Landau pole of the Yukawa coupling is above 30 TeV and that the η_Q production cross section times branching into diphotons (23) is 5 fb. The (red) cross-hatched tail is excluded because the Coulomb approximation fails. The (orange) region on the top of the figure is where the contribution to the decay width from the channel into diphotons is no longer negligible. The parameter space of the $Y_Q - N_{\text{TC}}$ plane is further constrained by the perturbativity condition $\alpha_{\text{TC}}(a_0^{-1}) < 1$ and by the requirement $\Gamma(\eta_Q \rightarrow ZH) < 10 \Gamma(\eta_Q \rightarrow \gamma\gamma)$. Note that within the allowed region of the parameter space we always have $\Gamma_{\text{total}} \leq 10 \text{ GeV}$ for each bound state. In fact, the constraint $\alpha_{\text{TC}}(a_0^{-1}) < 1$ itself requires $\Gamma_{\text{total}} < 11 \text{ GeV}$. The allowed region of the parameter space is significantly bounded; the hypercharge needs to lie between $2.26 \leq Y_Q \leq 2.53$ for $N_{\text{TC}} = 2$ and be within the range $2.17 \leq Y_Q \leq 2.31$ for $N_{\text{TC}} = 3$. Inside this region, the partial decay widths into W^+W^- and into ZZ are negligible compared to $\Gamma(\eta_Q \rightarrow \gamma\gamma)$ and hence the corresponding bounds are satisfied trivially.

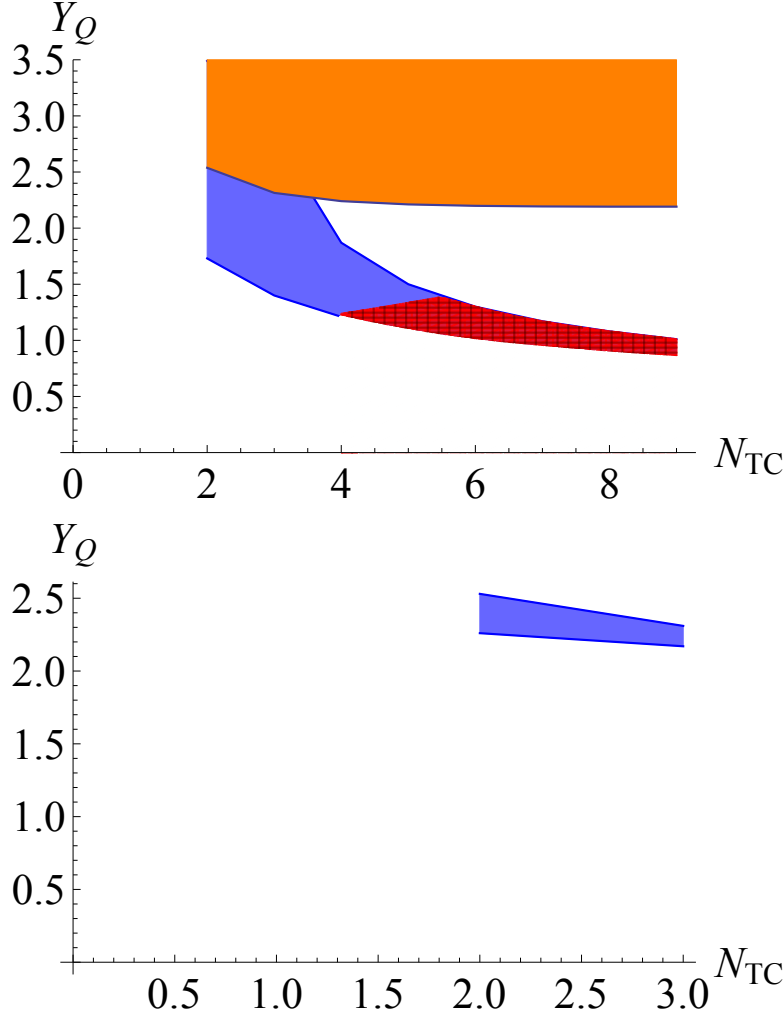


FIG. 1: **Top.** The (blue) banana-shape region is obtained by demanding that the Landau pole of the Yukawa coupling is above 30 TeV and that the production cross section of η_Q times its branching into diphotons is about 5 fb. The (red) cross-hatched tail is excluded because the Coulomb approximation fails. The (orange) band is where the contribution from decay into diphotons accounts for more than 10% of the total width. **Bottom.** Allowed region of the parameter space for $\Gamma(\eta_Q \rightarrow ZH) < 10 \Gamma(\eta_Q \rightarrow \gamma\gamma)$.

IV. CONSTRAINTS FROM COSMOLOGY

Primordial nucleosynthesis provides the earliest observationally verified landmark and constraints from big bang nucleosynthesis (BBN) can bound the parameters of beyond SM physics models [54]. Of particular interest here, the techni-gluons from the decay of Q -onium would hadronize to techni-glueballs G_{TC} of mass $M_G \sim 7\Lambda_{TC}$. We must then verify that the dominant techni-glueball decay, $G_{TC} \rightarrow \gamma\gamma$, does not drastically alter any of the light elemental abundances synthesized during BBN.

Note that if the G_{TC} decay takes place before BBN, then the photons injected into the plasma would rapidly redistribute their energy through scattering off background photons and through inverse Compton scattering. The thermalization process will be particularly

efficient at plasma temperatures above 1 MeV, which is the threshold for background e^+e^- pair annihilation, and which, incidentally, coincides with time of about 1 second. For $N_{\text{TC}} = 3$, the techni-glueball decay width is given by

$$\Gamma(G_{\text{TC}} \rightarrow \gamma\gamma) \approx \frac{(Y_Q + 1/2)^4 \alpha^2}{64\pi^3} \frac{M_G^3}{m_Q^2} \left(\frac{3M_G^3}{60m_Q^3} \right)^2; \quad (30)$$

the values for different N_{TC} are of the same order [55–57]. A straightforward substitution shows that in our model G_{TC} decay would not alter BBN as the techni-glueball lifetime $< 10^{-28}$ s.

V. LOOKING AHEAD

In this work we have attempted to associate the possible event excess in the diphoton invariant mass spectrum around 750 GeV, as indicated in LHC13 data, with bound states of a new asymptotically free gauge theory. In particular, we have constructed a minimal model with one generation of uncolored fermions Q , in which the Q fields transform as $SU(2)$ doublets and singlets, and are invariant under the $SU(N_{\text{TC}})$ transformation of a hidden strong gauge interaction to be explored during the LHC Run II data taking period. Since the new colorless fields transform as $SU(2)$ doublets they can acquire a mass through Yukawa interactions with the electroweak Higgs doublet, as quarks and leptons. We parametrized the Q -onium bound state using the Coulomb approximation and gave a numerical fit to the diphoton excess consistent with no strong coupling effects up to about 30 TeV. We have shown that allowed hypercharges lie within the range of $2.26 \leq Y_Q \leq 2.53$ for $N_{\text{TC}} = 2$ and $2.17 \leq Y_Q \leq 2.31$ for $N_{\text{TC}} = 3$. The smoking gun for the model would be the discovery of $\eta_Q \rightarrow ZH$.

In closing, we note that the Q Yukawa couplings drive the quartic Higgs coupling to negative values in the ultraviolet and the SM scalar effective potential develops an instability above about 100 TeV. As noted elsewhere [58] (see also [59, 60]) the potential instability of the electroweak vacuum can be evaded if the scalar sector contains a hidden heavy scalar singlet (with a large vacuum expectation value), which mixes with the SM Higgs doublet. The quartic interaction between the heavy scalar singlet and the Higgs doublet leads to a positive tree-level threshold correction for the Higgs quartic coupling, which is very effective in stabilizing the potential. In addition, the hidden scalar singlet could deliver mass terms for the vector-like hidden quarks Q and/or \bar{Q} [61].

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